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THE SINGLET g_2 STRUCTURE FUNCTION IN THE NEXT-TO-LEADING ORDER

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Abstract

Following a previous study of the one-loop factorization of the nonsinglet g_2 structure function of the nucleon, we present in this paper the next-to-leading order coefficient functions in the singlet sector. To obtain the result, the partonic processes of virtual Compton scattering off two and three "on-shell" gluons are calculated. A key step in achieving the correct factorization is to separate the correct twist-two contribution. The Burkardt-Cottingham sum rule is nominally satisfied at this order.

In lepton-nucleon deep-inelastic scattering with the electromagnetic current, two spin-dependent structure functions of the nucleon can be studied: $g_{1,2}(x_B, Q^2)$. $g_1(x_B, Q^2)$ is closely related to the spin structure of the nucleon and has been investigated extensively in the last decade. $g_2(x_B, Q^2)$ is present in processes involving transversely polarized nucleons and is a three-twist structure function in the sense that it contributes to physical observables at order $1/Q$ [1]. The experimental measurements of the g_2 structure function have been done by several collaborations [2]. There is a long history in debating the physics involved in $g_2(x_B, Q^2)$. As more results from quantum chromodynamics (QCD) are obtained and understood, it is generally accepted that $g_2(x_B, Q^2)$ probes quark and gluon correlations in the nucleon which cannot be accessed through Feynman-type incoherent parton scattering.

At leading-order, $g_2(x_B, Q^2)$ can be expressed in terms of a simple parton distribution $\Delta q_T(x)$,

$$g_T(x_B, Q^2) \equiv (g_1 + g_2)(x_B, Q^2) = \frac{1}{2} \sum_a e_a^2 \left(\Delta q_{aT}(x_B, Q^2) + \Delta \bar{q}_{aT}(x_B, Q^2) \right) \ , \tag{1}$$

where

$$\Delta q_T(x) = \frac{1}{2M} \int \frac{d\lambda}{2\pi} \mathrm{e}^{i\lambda x} \langle PS_\perp | \bar{\psi}(0) \gamma_\perp \gamma_5 \psi(\lambda n) | PS_\perp \rangle \ . \tag{2}$$

This result has lead to many incorrect interpretations of the g_2 physics. When the leading-logarithmic corrections are studied, it is found that $\Delta q_T(x, Q^2)$ mixes with other more complicated distributions under scale evolution [3]. In fact, $\Delta q_T(x, Q^2)$ is a special moment of a general class of parton distributions involving two light-cone variables [4, 5]. When the scale changes, only those general distributions evolve autonomously.

Thus a general factorization formula for g_2 is much more involved than the leading-order result shows. It should contain the generalized two-variable distributions, $K_i(x, y)$. Indeed, we shall write in general

$$g_T(x_B, Q^2) = \sum_i \int_{-1}^{1} \frac{dx}{x} \frac{dy}{y} \left\{ C_i \left(\frac{x_B}{x}, \frac{x_B}{y}; \alpha_s \right) K_i(x, y) + (x_B \to -x_B) \right\} \ , \tag{3}$$

where C_i are the coefficient functions. In a previous paper [6], we studied the factorization of the nonsinglet part at the one-loop order where there were two distributions, $K_{1a,2a}$, associated with each nonsinglet component a, and their one-loop coefficient functions were obtained for the first time.

In this paper, we study the one-loop factorization of the singlet part of $g_T(x_B, Q^2)$. The subject has previously been considered in Ref. [7], a comparison of the results will be made in the end of the paper. Contrary to the previous conclusion, the result of this paper represents a local operator product expansion. Throughout, we use the kinematics defined in [6]. We first define the singlet quark distributions,

$$K_{i\Sigma}(x, y) = \sum_a K_{ia}(x, y) \tag{4}$$

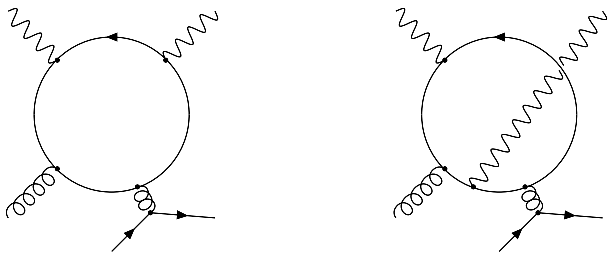

Figure 1: Feynman diagrams contributing to two quark and one gluon intermediate states. The intermediate longitudinal gluon is onshell, so a special propagator derived from the equation of motion has been employed. There are a total of six diagrams. In dimensional regularization, the sum is zero.

where the sum is over all quark flavors. Their one-loop coefficient functions are

$$C^{(0)}_{1,2\Sigma}\left(\frac{x_B}{x},\frac{x_B}{y}\right) = \bar{e}^2 y\delta(x-x_B) \ , \tag{5}$$

where $\bar{e}^2 = \sum_i e_i^2/N_f$ is the mean square quark charge and N_f is the number of active quark flavors. Using the relation

$$\Delta q_T(x) = \frac{2}{x}\int dy\,\{K_1(x,y)+K_2(x,y)\} \ , \tag{6}$$

it is easy to see that the result is consistent with Eq. (1).

The general strategy of higher-twist factorization beyond the leading orders has been presented in [6] and will not be repeated here. For the singlet factorization, we need to consider four classes of parton intermediate states: two quarks, two quarks-one gluon, two gluons, and three gluons. In the cases of two-quark and two-quark-one-gluon states, many of the relevant diagrams have been considered in [6]; we will not repeat that analysis here. In addition we must take into account the explicit singlet diagrams shown in Fig. 1. A detailed calculation shows that their contributions cancel. Thus the one-loop coefficient functions of $K_{i\Sigma}(x,y)$ remain the same as those of the non-singlet sector. In what follows, we focus entirely on two- and three-gluon intermediate states.

The diagrams involving two-gluon intermediate states are shown in Fig. 2. To isolate the $\mathcal{O}(1/Q)$ contribution, the incoming gluon is given a transverse momentum which is expanded to first order in the internal propagators. Our calculation shows that only the correlation function

$$\Gamma^\mu_{2gB}(x,y) = \delta(x-y)\int\frac{d\lambda}{2\pi}e^{i\lambda x}\langle PS|A^\mu_a(0)i\partial_\alpha A^\alpha_a(\lambda n)|PS\rangle \tag{7}$$

contributes, where α and μ take only transverse values.

The perturbative diagrams corresponding to three gluon intermediate states are shown in Fig. 3. The relevant correlation function is

$$\Gamma^\mu_{3gB}(x,y) = \int\frac{d\lambda}{2\pi}\frac{d\mu}{2\pi}e^{i\lambda x}e^{i\mu(y-x)}\langle PS|(-i)g_B f_{abc}A^\mu_a(0)A_{b\,\alpha}(\mu n)A^\alpha_c(\lambda n)|PS\rangle \ . \tag{8}$$

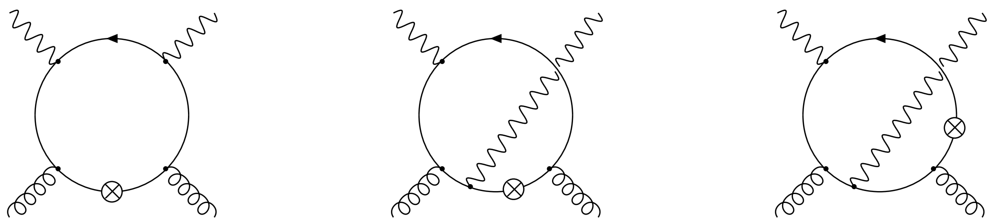

Figure 2: Two-gluon contributions to g_T. The cross $\otimes$ represents one transverse momentum operator insertion. There are a total of eight diagrams.

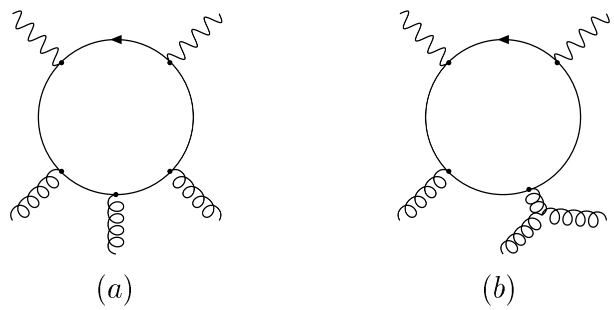

Figure 3: Feynman diagrams for three-gluon Compton scattering. There are a total of 24 diagrams represented in (a). The other 8, in figure (b), have the form of those in Fig. 1, and vanish for the same reason.

All fields and couplings in the above expressions are bare. Gauge invariance demands that the final result depends on the combination

$$\begin{aligned} \Gamma^\mu_{gB}(x,y) &= \Gamma^\mu_{2gB}(x,y) + \Gamma^\mu_{3gB}(x,y) \\ &= \int \frac{d\lambda}{2\pi}\frac{d\mu}{2\pi} \mathrm{e}^{i\lambda x}\mathrm{e}^{i\mu(y-x)} \langle PS|A^\mu_a(0) iD_{\alpha ab}(\mu n) A^\alpha_b(\lambda n)|PS\rangle \end{aligned} \tag{9}$$

of distributions, where $D^\mu_{\alpha ab} = \delta_{ab}\partial^\mu - gf_{acb}A_c$. In addition, gauge invariance demands that physical observables depend on moments of

$$\begin{aligned} xy\Gamma^\mu_{gB}(x,y) &= \int \frac{d\lambda}{2\pi}\frac{d\mu}{2\pi} \mathrm{e}^{i\lambda x}\mathrm{e}^{i\mu(y-x)} \langle PS|F^{+\mu}_a(0) iD_{\alpha ab}(\mu n) F^{+\alpha}_b(\lambda n)|PS\rangle \\ &= K_{gB}(x,y) i\epsilon^{\alpha\beta\gamma\mu} p_\alpha n_\beta S_\gamma \ . \end{aligned} \tag{10}$$

With appropriate insertions of light-cone gauge links, which can be generated by summing over states with additional longitudinally-polarized gluons, K_{gB} is gauge invariant. Our goal is to calculate its one-loop coefficient function.

The Compton amplitude for virtual photon scattering on a transversely polarized nucleon can be written as ($\epsilon^{0123} = +1$),

$$T^{\mu\nu} = -i\epsilon^{\mu\nu\alpha\beta} q_\alpha S_{\perp\beta} \frac{1}{\nu} S_T(x_B, Q^2) \ . \tag{11}$$

In QCD, we can write the singlet part as

$$S_T^\Sigma(x_B, Q^2) = \bar{e}^2 \int_{-1}^{1} \frac{dx}{x}\frac{dy}{y} \Bigg[\quad M_{1\Sigma}\left(\frac{x}{x_B}, \frac{y}{x_B}\right) K_{\Sigma 1B}(x,y) + M_{2\Sigma}\left(\frac{x}{x_B}, \frac{y}{x_B}\right) K_{\Sigma 2B}(x,y) + \frac{N_f}{y-x} M_g\left(\frac{x}{x_B}, \frac{y}{x_B}\right) K_{gB}(x,y)\Bigg] - (x_B \to -x_B) \ , \tag{12}$$

where the M's are perturbation series in α_s and have infrared poles. Once again, the tree and one-loop level expressions for $M_{i\Sigma}$ are the same as the nonsinglet case in Ref. [6].

The amplitude M_g starts at the one-loop level. To simplify the expression, we assume $|x_B| > 1$ so that the amplitude is purely real.

Since all Feynman diagrams are computed in bare perturbation theory, the result depends on the bare coupling g_B. We replace it with the renormalized coupling in the $\overline{\rm MS}$ scheme; the difference appears only at higher orders in $\alpha_s(\mu^2)$. A straightforward calculation yields

$$\begin{aligned} M_g^{(1)}(x,y) \;=\; & \frac{\alpha_s(\mu^2)T_F}{4\pi}\left(\frac{\bar\mu^2}{Q^2}\right)^{\epsilon/2}\frac{1}{2xy} \\ & \times \Bigg\{ \frac{2}{\epsilon}\Bigg[\left((y-2x)(1-x)+3y(y-x)-6\,\frac{y}{x}\,(y-x)\right)\log(1-x) \\ & + \left((4x-3y)(1-y)+4x(y-x)-8\,\frac{x}{y}\,(y-x)\right)\log(1-y) \\ & + \left(\frac{(x+y)(2x-3y)}{x-y}(1-(x-y))-xy\right)\log(1-(x-y))\Bigg] \\ & -3\left(y+2\,\frac{(y-x)^2}{x}\right)(1-x)\log(1-x) \\ & + \left(3(4x-3y)(1-y)+14x(y-x)-20\,\frac{x}{y}\,(y-x)\right)\log(1-y) \\ & + \frac{6x^2+xy-9y^2}{x-y}(1-(x-y))\log(1-(x-y)) \\ & - \frac{1}{2}\left((y-2x)(1-x)+3y(y-x)-6\,\frac{y}{x}\,(y-x)\right)\log^2(1-x) \\ & - \frac{1}{2}\left((4x-3y)(1-y)+4x(y-x)-8\,\frac{x}{y}\,(y-x)\right)\log^2(1-y) \\ & - \frac{1}{2}\left(\frac{(x+y)(2x-3y)}{x-y}(1-(x-y))-xy\right)\log^2(1-(x-y))\Bigg\} \ , \end{aligned} \tag{13}$$

where $\epsilon = 4-d$, $\bar\mu^2 = 4\pi e^{-\gamma_E}\mu^2$, and γ_E is the Euler constant. We have also introduced the generator normalization $T_F = 1/2$.

Now we want to show that $S_T^\Sigma(x_B, Q^2)$ is factorizable at the one-loop level, i.e., the infrared poles $1/\epsilon$ in M_g match the ultraviolet poles in K_{gB}. To this end, we use the infrared poles in M_g

to generate a scale evolution equation for the renormalized parton distributions,

$$
\begin{aligned}
\frac{d}{d\log\mu^2} & \int_{-1}^{1} dxdy \left(\frac{2}{x(x_B-x)}\right)\left(K_{1\Sigma}(x,y,\mu^2)+K_{2\Sigma}(x,y,\mu^2)\right) \\
= & \ \frac{\alpha_s(\mu^2)N_fT_F}{8\pi}\int_{-1}^{1}\frac{dxdy}{x^2y^2(x-y)} \\
& \times\Bigg[\left((y-2x)(x_B-x)+3y(y-x)-6\,\frac{x_By}{x}\,(y-x)\right)\log\left(1-\frac{x}{x_B}\right) \\
& \quad +\left((4x-3y)(x_B-y)+4x(y-x)-8\,\frac{x_Bx}{y}\,(y-x)\right)\log\left(1-\frac{y}{x_B}\right) \\
& +\left(\frac{(x+y)(2x-3y)}{x-y}\,(x_B-(x-y))-xy\right)\log\left(1-\frac{x-y}{x_B}\right)\Bigg]K_g(x,y)+\cdots\ , \qquad (14)
\end{aligned}
$$

where the ellipses denote the homogeneous part of the evolution. This contribution is given explicitly in Ref. [6]. In order to compare this result to the known twist-three evolution [4], we must first remove the twist-two contribution. Consider the twist-two operator

$$
\theta_{n+1}^{\mu_1\mu_2\cdots\mu_{n+1}} = F^{\alpha(\mu_1}iD^{\mu_2}\cdots iD^{\mu_n}i\tilde{F}^{\mu_{n+1})}_{\alpha}\ , \qquad (15)
$$

whose indices have been symmetrized and whose traces have been removed. Its matrix elements in the nucleon state are given by

$$
\langle PS|\theta_{n+1}^{\mu_1\cdots\mu_{n+1}}|PS\rangle = 2a_nP^{(\mu_1}\cdots P^{\mu_n}S^{\mu_{n+1})}\ . \qquad (16)
$$

It is the $\perp+\cdots+$ component of θ that appears in this transverse process :

$$
\begin{aligned}
\theta_{n+1}^{(i+\cdots)} = & \ \frac{2}{n+1}\epsilon^{ij}\Bigg[(n+1)iA^j(i\partial^+)^n i\partial^k A^k + iA^j(i\partial^+)^{n-1}D_\alpha F^{\alpha+} \\
& + f^{abc}A_a^j(i\partial^+)^{n-1}gA_b^k i\partial^+A_c^k + \sum_{m=0}^{n-1} f^{abc}A_a^j(i\partial^+)^m gA_b^k(i\partial^+)^{n-m}A_c^k\Bigg]\ . \qquad (17)
\end{aligned}
$$

Here, latin indices are understood to take values in the transverse dimensions. Expanding Eq. (14) in the large x_B limit, one arrives at evolution equations for the moments of the parton distributions. Removing the twist-two part of these equations via Eq. (17), one can obtain the evolution of the twist-three operators. A detailed check shows that our result is identical to that found in Ref. [4] obtained by studying the ultraviolet divergences present in the twist-three operators. We note here that this separation provides a new homogeneous term in the evolution of the singlet twist-three quark operators. In the absence of this term, the diagonal evolution of these operators would be identical to that in the nonsinglet sector since the contributions displayed in Fig. 1 vanish.

The final step of the calculation is to take the imaginary part of the factorized $S_T^{\Sigma}(x_B,Q^2)$ to get a factorized expression for the structure function $g_T^{\Sigma}(x_B,Q^2)$ in the physical region $x_B<1$.

We find the coefficient function

$$C_g^{(1)}\left(\frac{x_B}{x},\frac{x_B}{y}\right) = \frac{\alpha_s(\mu^2)T_F}{8\pi}\frac{1}{2xy}\left[3\left(y+2\,\frac{(y-x)^2}{x}\right)(x_B-x)\theta\left(\frac{x}{x_B}-1\right)\right.$$
$$-\left(3(4x-3y)(x_B-y)+14x(y-x)-20\,\frac{x_Bx}{y}\,(y-x)\right)\theta\left(\frac{y}{x_B}-1\right)$$
$$-\frac{6x^2+xy-9y^2}{x-y}(x_B-(x-y))\theta\left(\frac{x-y}{x_B}-1\right)$$
$$+\left((y-2x)(x_B-x)+3y(y-x)-6\,\frac{x_By}{x}\,(y-x)\right)\log\left(\frac{x}{x_B}-1\right)\theta\left(\frac{x}{x_B}-1\right)$$
$$+\left((4x-3y)(x_B-y)+4x(y-x)-8\,\frac{x_Bx}{y}\,(y-x)\right)\log\left(\frac{y}{x_B}-1\right)\theta\left(\frac{y}{x_B}-1\right)$$
$$\left.+\left(\frac{(x+y)(2x-3y)}{x-y}(x_B-(x-y))-xy\right)\log\left(\frac{x-y}{x_B}-1\right)\theta\left(\frac{x-y}{x_B}-1\right)\right]\ , \qquad (18)$$

where $\theta(x)$ is the step-function, which appears in Eq. 3 along with a factor of $1/(y-x)$.

To check the Burkhardt-Cottingham sum rule [8], we integrate $g_T^\Sigma(x_B,Q^2)$ over x_B. Assuming the integration over x and y can be interchanged with that of x_B, one obtains

$$\int_0^1 dx_B g_T^\Sigma(x_B,Q^2) = \frac{1}{2}\bar{e}^2\left(1-\frac{7}{2}C_F\frac{\alpha_s(Q^2)}{2\pi}\right)\langle PS|\sum_i\bar{\psi}_i\gamma_\perp\gamma_5\psi_i|PS\rangle\ , \qquad (19)$$

where γ_5 has been defined in the 't Hooft-Veltman scheme. The coefficient $7/2$ reduces to $3/2$ if we define γ_5 so that the nonsinglet axial current is conserved. Compared with the factorization formula for $g_1(x_B,Q^2)$ [9], we have the Burkhardt-Cottingham sum rule at one loop;

$$\int_0^1 dx_B g_2^\Sigma(x_B,Q^2) = 0\ . \qquad (20)$$

If the order of integration cannot be interchanged because of the singular behavior of the parton distributions at small x and y, the above sum rule may be violated. Indeed, some small x_B studies indicate such singular behavior [10].

Finally, we consider the next-to-leading order correction to the singlet part of the x^2 moment of $g_T(x_B,Q^2)$. In the leading order, it is well known:

$$\int_0^1 dx\ x^2 g_T^\Sigma(x,Q^2) = \frac{1}{3}\bar{e}^2\left(\frac{1}{2}a_{2\Sigma}(Q^2)+d_{2\Sigma}(Q^2)\right)\ , \qquad (21)$$

where a_2 is the second moment of the $g_1(x,Q^2)$ structure function and d_2 is a twist-three matrix element [11]. Using the coefficient functions found above, we have

$$\int_0^1 dx x^2 g_T^\Sigma(x,Q^2) = \frac{1}{3}\bar{e}^2\left\{\frac{a_{2\Sigma}(Q^2)}{2}\left[1+\frac{\alpha_s(Q^2)}{4\pi}\frac{7}{12}C_F\right]-\frac{a_{2g}(Q^2)}{2}\frac{\alpha_s(Q^2)}{4\pi}\frac{5}{3}N_fT_F\right.$$
$$\left.+d_{2\Sigma}(Q^2)\left[1+\frac{\alpha_s(Q^2)}{4\pi}\left(\frac{27}{4}C_A-\frac{29}{3}C_F+\frac{10}{3}N_fT_F\right)\right]\right\}\ . \qquad (22)$$

Using the next-to-leading result for $g_{1\Sigma}(x, Q^2)$ [9], we find

$$\int_0^1 dx x^2 \ \left(g_T^{\Sigma}(x, Q^2) - \frac{1}{3} g_1^{\Sigma}(x, Q^2) \right)$$
$$= \frac{1}{3} \bar{e}^2 d_{2\Sigma}(Q^2) \left\{ 1 + \frac{\alpha_s(Q^2)}{4\pi} \left(\frac{27}{4} C_A - \frac{29}{3} C_F + \frac{10}{3} N_f T_F \right) \right\} . \tag{23}$$

Notice that the combination of g_T and g_1 relevant to $d_{2\Sigma}$ receives no radiative correction, as in the nonsinglet sector. A detailed calculation shows that analogous results are valid for all higher moments as well [12]. This implies that the tree relation between $g_1(x, Q^2)$ and the twist-2 part of $g_2(x, Q^2)$ is respected at one-loop order. These results and their implications will be presented in a future communication [12].

The significance of the present result is as follows. In the leading order analysis of g_2, one just needs the leading-logarithmic evolution of $K_i(x, y)$ which is now well known [4, 5, 13], including its large N_c behavior [14]. In the next-to-leading order, one needs to know the coefficient functions and the two-loop evolution of $K_i(x, y)$. The former is now complete with Ref. [6] and the present paper. The latter has not yet been calculated, but in generally its effort is not as important as the coefficient function we calculate here.

Finally, let us add a few remarks on the comparison of our findings with a previous calculation of Ref. [7]. As has been noted earlier [15] the result of that paper is not complete since the contributions of twist-three two-gluon operators, i.e. the diagrams on Fig. 2, were not accounted for. Since the effect of two-gluon graphs on the final answer is reduced to a redefinition of the correlation function (schematically) from $\langle A^\perp_\mu A^\perp_\nu A^\perp_\rho \rangle$ to $\langle A^\perp_\mu D^\perp_\nu A^\perp_\rho \rangle$ without affecting the three-gluon coefficient function, we can directly compare both results. Before we do it, we observe that the basis of three-gluon functions used in [16, 7] is redundant, namely, in the decomposition

$$\langle A_\mu A_\nu A_\rho \rangle \propto N(x_1, x_2) g^\perp_{\mu\rho} \epsilon_{\nu-+\sigma} s^\perp_\sigma + \widetilde{N}(x_1, x_2) s^\perp_\rho \epsilon_{\mu\nu-+} + \dots , \tag{24}$$

obviously, the Lorentz structure in front of $\widetilde{N}$ can be expressed in terms of the first line using an identity

$$s^\perp_\rho \epsilon_{\mu\nu-+} = g^\perp_{\rho\nu} \epsilon_{\mu-+\sigma} s^\perp_\sigma - g^\perp_{\rho\mu} \epsilon_{\nu-+\sigma} s^\perp_\sigma . \tag{25}$$

Therefore, in final equations of [7] we have to put $\widetilde{N} = 0$ to be consistent with the present analysis. Moreover, the result in Ref. [16] must be also modified correspondingly.

Next, to have a correspondence with the correlation functions used here and in Ref. [7], we have identify the momentum fractions of the gluon lines as follows: $x = x_1$ and $y = x_1 - x_2$. For the latter convenience we introduce as well a new coefficient function

$$\mathcal{E}^{(1)}_g(x, y; x_B) \equiv \frac{1}{y - x} C^{(1)}_g \left(\frac{x_B}{x}, \frac{x_B}{y} \right) . \tag{26}$$

Since the function $N(x_1, x_2)$ is symmetric w.r.t. interchange of its arguments only symmetric part of $\mathcal{E}_g^{(1)}(x_1, x_1 - x_2; x_B)$ will survive. Thus, a symmetrization leads to the equation

$$\mathcal{E}_g^{(1)}(x_1, x_1 - x_2; x_B) + \mathcal{E}_g^{(1)}(x_2, x_2 - x_1; x_B) = \frac{\alpha_s}{8\pi} E_2(x_1, x_2; x_B), \tag{27}$$

where E_2 is the coefficient function from Ref. [7].

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